Leptogenesis during Axion Relaxation after Inflation

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In this talk, I present a novel and minimal alternative to thermal leptogenesis, which builds upon the assumption that the electroweak gauge bosons are coupled to an axion-like scalar field, as it is, for instance, the case in certain string compactifications. The motion of this axion-like field after the end of inflation generates an effective chemical potential for leptons and antileptons, which, in the presence of lepton number-violating scatterings mediated by heavy Majorana neutrinos, provides an opportunity for baryogenesis via leptogenesis. In contrast to thermal leptogenesis, the final baryon asymmetry turns out to be insensitive to the masses and CP-violating phases in the heavy neutrino sector. Moreover, the proposed scenario requires a reheating temperature of at least $\mathcal{O}(10^{12})$ GeV and it is, in particular, consistent with heavy neutrino masses close the scale of grand unification. This talk was given in February 2015 at *HPNP 2015* at Toyama University and is based on recent work (arXiv:1412.2043 [hep-ph]) in collaboration with A. Kusenko and T. T. Yanagida.

I. INTRODUCTION

The standard model (SM) fails to provide an explanation for the baryon-to-photon ratio in the present universe, $\eta_B^{\rm obs} \simeq 6 \times 10^{-10}$ [1], which serves as a major indication for new physics. Consequently, some new dynamical mechanism must be responsible for baryogenesis, i.e., the generation of a primordial baryon-antibaryon asymmetry in the early universe [2]. Most mechanisms proposed in the literature are devised so as to satisfy the three famous Sakharov conditions [3]: (i) violation of baryon number B (or lepton number L in the case of leptogenesis), (ii) violation of C as well as of CP invariance, and (iii) departure from thermal equilibrium. As it turns out, it is, however, not mandatory to fulfill these three conditions in order to successfully generate a baryon asymmetry. The point is that Sakharov's conditions are based on the assumption of CPT invariance, which means that one is actually able to circumvent them in case CPT is spontaneously broken. This idea has been pioneered by Cohen and Kaplan in their scenario of *spontaneous baryogenesis* [4], where the baryon asymmetry is generated in thermal equilibrium; and since then, it has been studied and expanded upon by many authors [5–8]. For example, Kusenko et al. have recently shown how the CPT violation during the phase of SM Higgs relaxation after the end of inflation can be used for the realization of baryogenesis via leptogenesis [9]. In this talk, I will draw upon this earlier work and demonstrate that it can be easily generalized to the case of generic axion-like scalar fields relaxing from large initial field values in the course of reheating; further details pertaining to our analysis can be found in our recent paper [10] as well as in another forthcoming publication.

In an expanding universe at nonzero temperature, CPT invariance can be easily broken spontaneously by introducing a pseudoscalar field, $a(t,\vec{x})$, which couples derivatively to the fermion current j^μ in the Lagrangian,

$$\mathcal{L} \supset \frac{1}{f_a}\,\partial_\mu a\, j^\mu\,, \quad j^\mu = \sum_f \bar{\psi}_f \gamma^\mu \psi_f\,, \tag{1}$$

with f_a being some cut-off scale. Imposing spatial homogeneity, $a = a(t)$, and assuming that the classical background is given cause to evolve with nonzero velocity, $\dot{a} \neq 0$, (which is readily done in the early universe, as we will review shortly) this coupling turns into an effective chemical potential $\mu_{\rm eff}$ for the fermion number,

$$\mathcal{L} \supset \frac{1}{f_a}\,\dot{a}\, j^0 = \mu_{\rm eff}\, n_F\,, \quad \mu_{\rm eff} = \frac{\dot{a}}{f_a}\,, \quad j^0 \equiv n_F = n_f - n_{\bar{f}}\,, \tag{2}$$

which shifts the energy levels of fermions f and antifermions $\bar{f}$ w.r.t. each other. In thermal equilibrium, the minimum of the free energy is therefore obtained for a nonzero fermion-antifermion asymmetry n_F,

$$n_{f,\bar{f}}^{\rm eq} \sim T^3\left(1 \pm \frac{\mu_{\rm eff}}{T}\right)\,, \quad n_F^{\rm eq} = n_f^{\rm eq} - n_{\bar{f}}^{\rm eq} \sim \mu_{\rm eff}\, T^2\,. \tag{3}$$

As observed by Cohen and Kaplan, this result may serve as a basis for the successful generation of the baryon asymmetry. However, in order to arrive at a realistic model, one first of all has to address three important questions: (i) what is the nature of the field a and the origin of the derivative coupling in Eq. (1), (ii) how is the field a set in motion, and (iii) what kind of interactions drive the number density n_F towards its equilibrium value $n_F^{\rm eq}$? In the following, I shall discuss each of these issues in turn, cf. Sec. II, which will eventually lead us to an interesting alternative to thermal leptogenesis [11]. In Sec. III, I will then sketch the parameter dependence of the final baryon asymmetry in our model; and in Sec. IV, I will conclude and give a brief outlook.

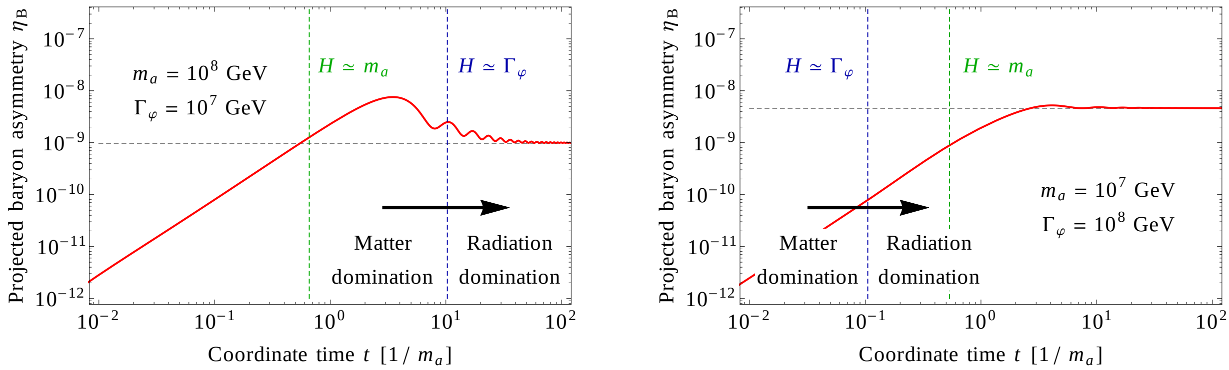

FIG. 1: Evolution of the (instantaneous) baryon asymmetry (projected onto its would-be present-day value) as a function of time for $m_a \gg \Gamma_\varphi$ **(left panel)** and $m_a \ll \Gamma_\varphi$ **(right panel)**. Here, $f_a = 3 \times 10^{14}\,\text{GeV}$ in both panels.

II. NOVEL AXION-DRIVEN LEPTOGENESIS MECHANISM

Let us integrate the derivative interaction in Eq. (1) by parts, $\mathcal{L} \supset 1/f_a\,\partial_\mu a\, j^\mu \rightarrow -a/f_a\,\partial_\mu j^\mu$. This illustrates that the CPT-violating coupling required for baryogenesis is equivalent to a coupling of the pseudoscalar a to the divergence of the fermion current j^μ. The field a is thus naturally identified as an axion-like field, or simply "axion" [12], which couples to the anomaly of the fermion number $F = 3\,B + L$. From this point of view, the cut-off scale f_a in Eq. (1) is immediately recognized as the decay constant of the axion field a. Moreover, we note that the electroweak anomalies of the baryon number $U(1)_B$ and lepton number $U(1)_L$ in the SM allow us to recast the axion coupling to $\partial_\mu j^\mu$ as a coupling to the electroweak field strength tensors $W_{\mu\nu}$ and $B_{\mu\nu}$,

$$\mathcal{L} \supset -\frac{a}{f_a}\,\partial_\mu j^\mu \rightarrow \frac{a}{f_a}\frac{N_f}{8\pi^2}\left(g_2^2\,W_{\mu\nu}\tilde{W}^{\mu\nu} - g_1^2\,B_{\mu\nu}\tilde{B}^{\mu\nu}\right)\,, \tag{4}$$

where $N_f = 3$ is the number of SM fermion generations and with g_2 and g_1 denoting the electroweak gauge couplings. Interactions of this form may, for instance, arise in string theory, which always features at least one (model-independent) axion [13] associated with the Green-Schwarz mechanism of anomaly cancellation [14]. This axion couples to all gauge groups with universal strength $f_a \sim 10^{16}\,\text{GeV}$ [15]. Besides that, string theory may also give rise to a multitude of further axions fields coupling to different gauge groups with nonuniversal strength [13, 16]. The couplings of these model-dependent axions are then determined by the gauge structure as well as the details of the compactification. For our purposes, the upshot of these considerations is that a certain linear combination of stringy axions may very well end up coupling to $F\tilde{F}$, where $F = W, B$. In the following, we shall therefore identify the field a in Eq. (1) with just this linear combination and take the above string-based argument to be the origin of the coupling $\mathcal{L} \supset a/f_a\,F\tilde{F} \leftrightarrow 1/f_a\,\partial_\mu a\, j^\mu \leftrightarrow \mu_{\text{eff}}\, n_F$ in the Lagrangian.

The dynamics of the axion background in the early universe are governed by its classical equation of motion,

$$\ddot{a} + 3\,H\,\dot{a} = -\partial_a V_{\text{eff}}(a)\,, \quad V_{\text{eff}} \simeq \frac{1}{2}m_a^2 a^2\,, \tag{5}$$

where the effective potential V_{eff} and hence the axion mass $m_a \simeq \Lambda_H^2/f_a$ may, for instance, originate from instanton effects in a strongly coupled hidden sector featuring a dynamical scale Λ_H. Assuming that the PQ-like symmetry associated with the flat axion direction is broken sufficiently early before the end of inflation (and not restored afterwards), the initial axion field value $a_0 = (0 \cdots 2\pi)f_a$ becomes constant on superhorizon scales. For definiteness, we shall therefore take a_0/f_a to be 1 in the entire observable universe at the end of inflation. As we will see shortly, the baryon asymmetry produced during reheating is going to depend on a_0. Because of that, we have to ensure that the baryonic isocurvature perturbations induced by the quantum fluctuations δa of the axion field around its homogeneous background a_0 remain below the observational bound [17]. This constrains the Hubble rate H_{inf} during inflation: $H_{\text{inf}}/(2\pi)/a_0 \lesssim 10^{-5}$ or equivalently $H_{\text{inf}} \lesssim 6\times 10^{11}\,\text{GeV}\left(f_a/10^{15}\,\text{GeV}\right)$. At the same time, we have to require that $m_a \lesssim H_{\text{inf}}$, so that during inflation the Hubble friction term on the left-hand side of Eq. (5) outweighs the potential gradient on the right-hand side of Eq. (5). After the end of inflation, the Hubble rate then begins to drop, until, around $H \simeq m_a$, the axion begins to coherently oscillate around the minimum of its effective potential, $a = 0$, with frequency $\omega = m_a$. During this stage of *axion*

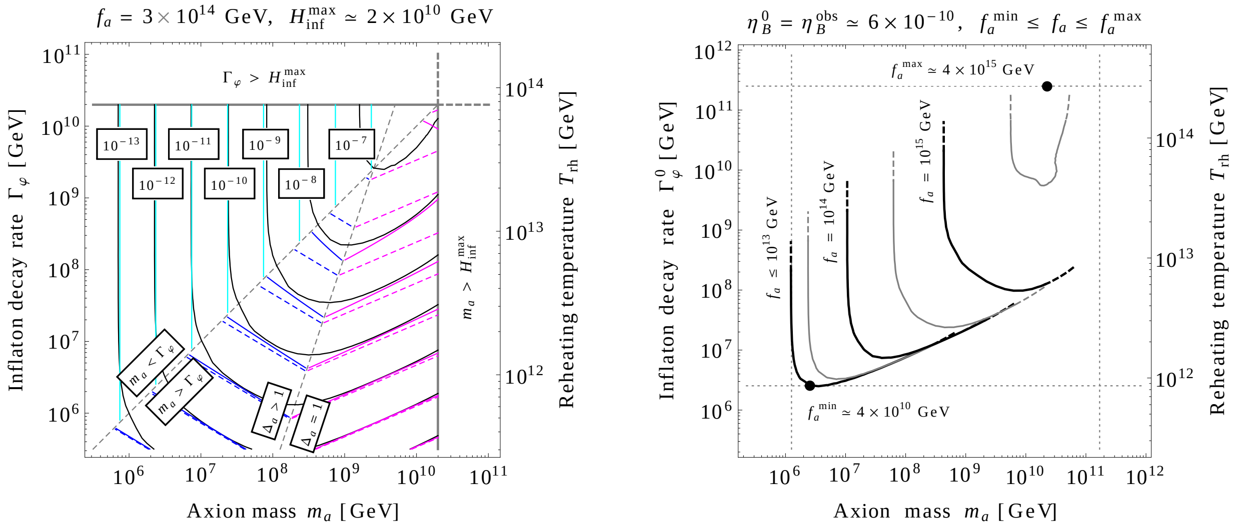

FIG. 2: **(Left panel)** Contour plot of the final baryon asymmetry η_B^0. The black (bent) contours represent the full numerical result, while the colorful (straight) contours depict the analytical estimate according to Eqs. (9), (10) and (11). The effect of washout is illustrated by the difference between the dashed ($\kappa = 0$) and solid ($\kappa \neq 0$) lines. **(Right panel)** Contour lines of successful leptogenesis ($\eta_B^0 = \eta_B^{\rm obs}$) for different values of f_a. The dashed segments along the individual contours indicate where either m_a or Γ_φ becomes comparable to the maximally allowed Hubble rate, $H_{\rm inf}^{\rm max} \simeq 2\pi\, 10^{-5}\, f_a$.

relaxation, the axion field therefore evolves with nonzero velocity in its potential, which temporarily induces an effective chemical potential for the fermion number, as anticipated at the beginning of this talk, cf. Eq. (2).

In order to make use of $\mu_{\rm eff} = \dot{a}/f_a$ for the purposes of baryogenesis, it is important that there be an *external source* of baryon or lepton number violation that proceeds at a much faster rate Γ than the motion of the axion field in its effective potential, $\Gamma \gg \dot{a}/a$. Only then do the axion oscillations act as an adiabatic background, so that $\dot{a}/f_a$ can be interpreted as an effective chemical potential [7]. Here, a minimal choice to satisfy this requirement is to rely on L violation through the s- and t-channel exchange of heavy Majorana neutrinos N_i,

$$\Delta L = 2: \quad \ell_i \ell_j \leftrightarrow N_k^* \leftrightarrow HH\,, \; \ell_i H \leftrightarrow N_k^* \leftrightarrow \bar{\ell}_j \bar{H}\,, \quad \ell_i^T = \begin{pmatrix} \nu_i & e_i \end{pmatrix}\,, \; H^T = \begin{pmatrix} h_+ & h_0 \end{pmatrix}\,, \; i,j,k = 1,2,3\,. \quad (6)$$

These processes are guaranteed to be present in the bath as long as one believes in the seesaw mechanism as the correct explanation for the small neutrino masses in the SM [18]. In order to separate the leptogenesis mechanism under study form the contributions from ordinary thermal leptogenesis, we shall assume that all right-handed neutrinos N_i acquire Majorana masses M_i close to the scale of grand unification (GUT), $M_i \sim \mathcal{O}\,(0.1 \cdots 1)\,\Lambda_{\rm GUT} \sim 10^{15} \cdots 10^{16}\,{\rm GeV}$, so that none of them is actually ever thermally produced. For center-of-mass energies $\sqrt{s} \ll M_i$, the thermally averaged cross section of the $\Delta L = 2$ lepton-Higgs scattering processes in Eq. (6), $\sigma_{\rm eff} \equiv \langle \sigma_{\Delta L = 2}\, v \rangle$, is then practically fixed [19] by the experimental data on the SM neutrino sector [20],

$$\sigma_{\rm eff} \approx \frac{3}{32\pi} \frac{\bar{m}^2}{v_{\rm ew}^4} \simeq 1 \times 10^{-31}\,{\rm GeV}^{-2}\,, \quad \bar{m}^2 = \sum_{i=1}^{3} m_i^2 \approx \Delta m_{\rm atm}^2 \simeq 2.4 \times 10^{-3}\,{\rm eV}^2\,, \quad v_{\rm ew} \simeq 174\,{\rm GeV}\,. \quad (7)$$

Correspondingly, the evolution of the L number density n_L is described by the following Boltzmann equation,

$$\dot{n}_L + 3\,H\,n_L \simeq -\Gamma_L \left(n_L - n_L^{\rm eq}\right)\,, \quad \Gamma_L = 4\,n_\ell^{\rm eq} \sigma_{\rm eff}\,, \quad n_\ell^{\rm eq} = \frac{2}{\pi^2} T^3\,, \quad n_L^{\rm eq} = \frac{4}{\pi^2} \mu_{\rm eff} T^2\,, \quad (8)$$

where we have approximated the lepton and L number densities in thermal equilibrium, $n_\ell^{\rm eq}$ and $n_L^{\rm eq}$, by their corresponding expressions in the classical Boltzmann approximation. Notice that the production term on the right-hand side of this equation, $\Gamma_L n_L^{\rm eq} \propto \sigma_{\rm eff}\, \mu_{\rm eff}\, T^5$, is largely independent of the details of the neutrino sector. It does, in particular, not depend on the amount of CP violation in the neutrino sector nor on the heavy neutrino mass spectrum. At the same time, it increases linearly with the light neutrino mass scale $\bar{m}$. The usual bound on this mass scale from thermal leptogenesis, $\bar{m} \lesssim 0.2\,{\rm eV}$, (where it ensures that dangerous washout processes do not become too strong [21]) hence does not apply in our scenario. The absolute neutrino mass scale will soon be probed experimentally (on earth [22] as well as in the sky [23]). This, therefore, entails the intriguing possibility to test our model against the conventional scenario of thermal leptogenesis in the near future.

III. PARAMETER DEPENDENCE OF THE FINAL BARYON ASYMMETRY

Subsequent to its generation (according to Eq. (8)), the lepton asymmetry is partly converted into a baryon asymmetry by means of electroweak sphalerons. The final present-day baryon asymmetry η_B^0 is then given as

$$\eta_B^0 = \frac{n_B^0}{n_\gamma^0} = c_{\rm sph}\,\frac{g_{*,s}^0}{g_*}\,\eta_a^L \simeq 0.013\,\eta_L^a\,, \quad c_{\rm sph} = \frac{28}{79}\,, \quad g_{*,s}^0 = \frac{43}{11}\,, \quad g_* = \frac{427}{4}\,, \tag{9}$$

where $c_{\rm sph}$, g_* and $g_{*,s}^0$ denote the SM sphaleron conversion factor as well as the effective number of relativistic degrees of freedom at high and low temperatures, respectively. Moreover, η_L^a represents the final lepton asymmetry after the late-time decay of the axion field at times around $t \sim \Gamma_a^{-1}$, where $\Gamma_a \simeq g_2^2/(256\,\pi^4)\,m_a^3/f_a^2$. In general, η_L^a does not correspond to the equilibrium number density at the same time, $\eta_L^{\rm eq}$, as the efficiency of leptogenesis typically begins to cease before the equilibrium value is actually reached, so that $\eta_a^L \ll \eta_a^{\rm eq}$.

For fixed values of a_0 and f_a, the final lepton asymmetry ends up depending on two parameters: the axion mass m_a as well as the reheating temperature $T_{\rm rh}$, which is, in turn, determined by the inflaton decay rate, $T_{\rm rh} \simeq 0.3\,\sqrt{\Gamma_\varphi M_{\rm Pl}}$. We infer the precise parameter dependence of η_L^a by numerically solving Eqs. (5) and (8) together with the Friedmann equation for the scale factor as well as the Boltzmann equations for the number densities of inflaton particles and relativistic SM particles, respectively. As it turns out, our exact numerical result can be very well reproduced by the following analytical expression, cf. also the left panel of Fig. 2,

$$\eta_L^a = C\,\Delta_a^{-1}\Delta_\varphi^{-1}\,\eta_L^{\rm max}\,e^{-\kappa}\,, \quad \eta_L^{\rm max} \simeq \frac{\sigma_{\rm eff}}{g_*^{1/2}}\frac{a_0}{f_a}\,m_a\,M_{\rm Pl} \times \min\left\{1,\,(\Gamma_\varphi/m_a)^{1/2}\right\}\,, \tag{10}$$

with C being a numerical fudge factor of $\mathcal{O}(1)$. $\eta_L^{\rm max}$ denotes the all-time maximum of the lepton asymmetry, which is reached around the time when the axion oscillations set it, i.e., at $t \sim t_{\rm osc} \simeq m_a^{-1}$. Note that it is rather insensitive to both a_0 and f_a, as it only depends on the ratio a_0/f_a, which is expected to be of $\mathcal{O}(1)$. Furthermore, Δ_φ and Δ_a account for the dilution of $\eta_L^{\rm max}$ in the course of inflaton and axion decays, respectively,

$$\Delta_\varphi \simeq \max\left\{1,\,(m_a/\Gamma_\varphi)^{5/4}\right\}\,, \quad \Delta_a \simeq \max\left\{1,\,\frac{8\pi^3}{g_2^2}\frac{f_a\,a_0^2}{m_a\,M_{\rm Pl}^2} \times \min\left\{1,\,(\Gamma_\varphi/m_a)^{1/2}\right\}\right\}\,. \tag{11}$$

Here, Δ_φ reflects the interplay between leptogenesis and reheating, cf. Fig. 1. For $m_a \gtrsim \Gamma_\varphi$, the axion begins to oscillate before the end of reheating and the initial asymmetry becomes diluted due to the entropy production in inflaton decays. For $m_a \lesssim \Gamma_\varphi$, on the other hand, the axion oscillations only set in after the end of reheating and the final asymmetry becomes independent of the inflaton decay rate. Meanwhile, Δ_a begins to play a role for f_a values around 3×10^{13} GeV, cf. the right panel of Fig. 2. For smaller values of f_a, we always have $\Delta_a = 1$ in the entire parameter region of interest. The factor $e^{-\kappa}$, finally, accounts for the efficiency of the washout term, $-\Gamma_L n_L$, on the right-hand side of Eq. (8). For $m_a \gtrsim \Gamma_\varphi$, κ can be roughly estimated as $\kappa \sim T_{\rm rh}/T_L$, where $T_L \sim 1/\left(\sigma_{\rm eff} M_{\rm Pl}\right) \sim 10^{13}\,$GeV is the typical temperature scale of leptogenesis, while, for $m_a \lesssim \Gamma_\varphi$, we have $\kappa \sim 1$. A better analytical understanding of washout in our scenario is, however, still pending.

Successful leptogenesis restricts the axion decay constant f_a to take a value within the following range,

$$4 \times 10^{10}\,\text{GeV} \lesssim f_a \lesssim 4 \times 10^{15}\,\text{GeV}\,. \tag{12}$$

which translates into allowed ranges for m_a, Γ_φ and $T_{\rm rh}$, cf. the right panel of Fig. 2, which are all very well consistent with typical string axion models. We note that, for smaller values of f_a, it is not possible to generate a sufficiently large baryon asymmetry, while keeping the baryonic isocurvature perturbations small enough. Likewise, for larger values of f_a, the dilution of the asymmetry in the late-time decay of the axion is too strong.

IV. CONCLUSIONS AND OUTLOOK

While thermal leptogenesis typically operates at $T_{\rm rh} \sim 10^{10}\,$GeV, the requirement of a large rate of L violation, $\Gamma_L \gg H$, pushes $T_{\rm rh}$ to values at least of $\mathcal{O}\left(10^{12}\right)$ GeV in our scenario. Furthermore, our final baryon asymmetry turns out be independent of the amount of CP violation in the neutrino sector as well as of the N_i mass spectrum. On top of that, the usual bound on $\bar{m}$ from thermal leptogenesis does not apply in our case. The presented model should therefore be regarded as an attractive alternative to thermal leptogenesis in case the latter should begin to look less favorable from the experimental point of view! Beyond that, further work is needed: it remains, for instance, to be seen how the required high $T_{\rm rh}$ could be possibly accommodated in a supersymmetric version of our model. A further intriguing question, which we are currently investigating, is whether the role of the axion field a could not be equally played by the inflaton. This would result in an even more minimal scenario.

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